in SUGRA Models with Nonuniversal Gaugino Masses

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Abstract

The constraints of electric dipole moments (EDMs) of electron and neutron on the parameter space in supergravity (SUGRA) models with nonuniversal gaugino masses are analyzed. It is shown that with a light sparticle spectrum , the sufficient cancellations in the calculation of EDMs can happen for all phases being order of one in the small tanβ case and all phases but ϕ_μ ($|\phi_\mu| \lesssim \pi/6$) order of one in the large tanβ case. This is in contrast to the case of mSUGRA in which in the parameter space where cancellations among various SUSY contributions to EDMs happen $|\phi_\mu|$ must be less than $\pi/10$ for small $tan\beta$ and $\mathcal{O}(10^{-2})$ for large $tan\beta$. Direct CP asymmetries and the CP violating normal polarization of lepton in $B \to X_s l^+ l^-$ are investigated in the models. In the large tanβ case, A^2_{CP} and P_N for l=μ (τ) can be enhanced by about a factor of ten (ten) and ten (three) respectively compared to those of mSUGRA.

Recent observation of Re(ϵ'/ϵ) by KTev collaboration [1] definitely confirms the earlier NA31 experiment[2]. This direct CP violation measurement in the Kaon system can be accommodated by the CKM phase in standard model (SM) within the theoretical uncertainties. However, the CKM phase is not enough to explain the matter-antimatter asymmetry in the universe and gives the contribution to EDMs much smaller than the limits of EDMs of electron and neutron. One needs to have new sources of CP violation and examine their phenomenological effects.

There exist new sources of CP violation in SUSY theories which come from the phases of soft SUSY breaking parameters. It is well-known for a long time that in order to satisfy the current experimental limits on EDMs of electron and neutron SUSY CP-violating phases have to be much smaller($\lesssim 10^{-2}$) unless sfermion masses of the first and second generations are very large (> 1 Tev) [3]. Recently it has been shown that various contributions to EDMs cancel with each other in significant regions of the parameter space so that the current experimental limits on the EDM of electron (EDME) [4] and the EDM of neutron (EDMN) [5],

$$|d_e| < 4.3 \times 10^{-27} ecm \tag{1}$$

and

$$|d_n| < 6.3 \times 10^{-26} ecm, \tag{2}$$

can be satisfied for SUSY models with SUSY phases of order one and relatively light sparticle spectrum (< 1 Tev) [6, 7]. In mSUGRA even in the parameter space where cancellation among various SUSY contributions to neutron EDM(EDMN) happens $|\phi_\mu|$ must be less than $\pi/10$ for small $tan\beta$[6] and $\mathcal{O}(10^{-2})$ for large $tan\beta$ [8] while the allowed range of ϕ_{A_0} is almost unconstrained. Brhlik et al. pointed out that more sufficient cancellations happen in MSSM if gaugino masses are complex [9]. In the letter we consider cancellation phenomena in SUGRA with nonuniversal gaugino masses.

CP violation has so far only been observed in K system. It is one of the goals of the B factories presently under construction to discover and examine CP violation in the B system. Direct CP violation in $B \to X_s l^+ l^-$ in SM has been examined and the result is that it is unobservablly small [10]. In mSUGRA the CP assymmetry of branching ratio on $B \to X_s l^+ l^-$has been given in [22]. A detailed analysis of SUSY contributions to CP Violation in semileptonic B decays has been performed using the mass insertion approximation in [23]. Direct CP asymmetries and the CP violating normal polarization of lepton in $B \to X_s l^+ l^-$ in mSUGRA with CP-violating phases are investigated in our previous paper [8]. In the letter we extend the investigation to SUGRA models with nonuniversal gaugino masses after studying the allowed regions of the parameter space in the models by EDM data.

In order to concentrate on the effects of the phases arising from complex gaugino masses we limit ourself to a class of SUGRA models with nonuniversal gaugino masses in which gaugino masses at high energy scale (GUT scale) are nonuniversal but scalar masses and trilinear couplings at GUT scale are still universal. Such a class of effective SUGRA models can naturally arise from string models [11]. In this class of models, compared to the mSUGRA, there are two more new independent phases [9] which we choose to be ϕ_1 and ϕ_3, the phases of gaugino masses M_1 and M_3, in addition to the phases ϕ_μ and ϕ_{A_0}. From the one loop renormalization group equations (RGEs) for M_i (i=1,2,3) [12]

$$\frac{dM_i}{dt} = \frac{1}{4\pi} b_i \alpha_i M_i \qquad i = 1, 2, 3 \tag{3}$$

where $\alpha_i = \frac{g_i^2}{4\pi}$, $t = ln(Q^2/M_{GUT}^2)$, we know that the phases of M_i do not run, like the phase of μ.

Let us recall the cancellation mechanism for EDME. There are only two contributions, the chargino (-sfermion loop) and neutrilino (-sfermion loop) contributions, to the EDME. The chargino contribution involves gaugino -Higgsino(g-h) mixing. The neutrilino contribution involves not only gaugino-Higgsino mixing but also gaugino-gaugino(g-g) mixing. The chargino contribution and the part of the neutrilino contribution which comes from g-h mixing have automatically opposite sign because of the opposite sign of μ in chargino and neutrilino mass matrices. In general, the chargino contribution in magnitude is significantly larger than the part of neutrilino contribution. Therefore, as pointed out in ref. [9], a cancellation can happen only if the another part of the neutrilino contribution which comes from the g-g mixing can balance some of the difference between the two contributions. For EDME the neutrilino contribution which comes from the g-g mixing is proportional to [9]

$$\frac{1}{m_{\tilde{e}}^2} m_e [A_e sin(\phi_{A_e} - \phi_1) + |\mu| tan\beta sin(\phi_\mu + \phi_1)] \tag{4}$$

Therefore, given ϕ_μ, the sign of the contribution can be controlled by choosing ϕ_1 and ϕ_{A_0} and the magnitude of the contribution can increase by increasing μ and/or A_e. Because the chargino contribution is dependent on μ and independent on ϕ_{A_e}, it is sufficient to have a cancellation that the magnitude of the A_e term (first term) in eq.(4) is comparable to that of the $\mu tan\beta$ term (second term) in eq.(4). This is easy to be down in MSSM in which A_e and μ are free parameters. Thus, an almost exact cancellation can occur for the whole range of ϕ_μ. That is exactly what happens in MSSM [9].

However, in SUGRA models low energy properties are determined by running RGEs from the high energy scale to the electroweak(EW) scale and the radiative breaking mechanism of the EW symmetry puts constraints on CP- violating phases. As long as we limit our discussion to mass spectra less than than 1 Tev, M_3 and A_0(hence A_e) can not be too large. For small tan β(say $\lesssim$ 2), the sufficient condition (i.e., the two terms in eq.(4) have size of the same order) can easily be realized in the almost whole range of ϕ_μ by choosing ϕ_{A_0} and ϕ_1. For moderate and large tanβ, it is difficult for the condition to be realized due to the limited values of A_0 (hence A_e) so that only for some limited ranges of ϕ_μ the EDM constraint can be satisfied. The similar (but more complicated) situation occurs for EDMN with appropriate ϕ_3 as well as ϕ_{A_0} chosen.

In order to show the important role of ϕ_{A_0} played in the cancellation mechanism, in fig.1a and 1b we display EDME as function of ϕ_1 for ϕ_{A_0}=0 and different ϕ_μ for both small tanβ (2) and large tanβ (30) cases. We can see from the fig.1 that most of the range of ϕ_μ is excluded by EDME in both cases. For EDMN as function of ϕ_3, similar results are obtained. That is, like ϕ_1, for positive ϕ_μ cancellations happen in some narrow ranges within [π,2π] of ϕ_3 and within [0,π] for negative ϕ_μ. When we vary the values of ϕ_{A_0} we achieve the above mentioned results: almost whole range of ϕ_μ is allowed by EDME and EDMN for small tanβ and $|\phi_\mu| \lesssim \pi/6$ for large tanβ(see fig.2). Moreover, because $\phi_1(\phi_3)$ is correlated with ϕ_μ, we find that with varying ϕ_μ the whole range of ϕ_1 and ϕ_3 can be allowed by EDM constraints. For large tanβ (30) case largest $|\phi_\mu|$(about $\pi/6$) correspond to ϕ_1 and ϕ_3 around $\mp\pi/2\mp\pi/6$, while ϕ_1 and ϕ_3 are around $\pm\pi/2$ when ϕ_μ about ∓ 0.4, and when ϕ_μ is about ± 0.2 they are around $\mp\pi/4$. The correlated values of ϕ_3 and ϕ_μ are needed in analyses of $B \to X_s l^+ l^-$ and $B \to X_s \gamma$ below. Correlation between ϕ_μ and tanβ, with the absolute value of soft breaking terms chosen as those in fig1 and appropriate phases chosen, is shown in fig.2 where all of the points are allowed by the experimental bounds on EDME and EDMN. One can see in the figure that ϕ_μ becomes more constrained as $tan\beta$ is increased. Nevertheless, for $tan\beta$ larger than 6 the allowed regions of ϕ_μ are almost unchanged, which means that effects of the A_e term (and A_d term in the case of EDMN) in eq.(4) are relatively small and the balance is provided by the $\mu tan\beta$ term in eq.(4). Since we also consider the large $tan\beta$ case, we include the two loop contribution given by D. Chang et.al[13]. But the numerical calculations in the regions of the parameter space in which one loop EDMs satisfy the current experimental limits due to the cancellation mechanism show that it is very small compared to one loop contributions.

We now turn to the calculations of the CP violation in $B \to X_s l^+ l^-$. The direct CP asymmetries in decay rate and backward-forward assymmetry for $B \to X_s l^+ l^-$ and $\overline{B} \to \overline{X}_s l^+ l^-$ are defined by [8, 14]

$$A^1_{CP}\ (\hat{s}) = \frac{\mathrm{d}\Gamma/\mathrm{d}\hat{s} - \mathrm{d}\overline{\Gamma}/\mathrm{d}\hat{s}}{\mathrm{d}\Gamma/\mathrm{d}\hat{s} + \mathrm{d}\overline{\Gamma}/\mathrm{d}\hat{s}} = \frac{D(\hat{s}) - \overline{D}(\hat{s})}{D(\hat{s}) + \overline{D}(\hat{s})},$$

$$A^2_{CP}\ (\hat{s}) = \frac{A(\hat{s}) - \overline{A}(\hat{s})}{A(\hat{s}) + \overline{A}(\hat{s})} \tag{5}$$

where

$$\begin{aligned} A(\hat{s}) &= 3\sqrt{1 - \frac{4t^2}{\hat{s}}}\frac{E(\hat{s})}{D(\hat{s})}, \\ D(\hat{s}) &= 4|C_7|^2(1 + \frac{2}{\hat{s}})(1 + \frac{2t^2}{\hat{s}}) + |C_8^{eff}|^2(2\hat{s} + 1)(1 + \frac{2t^2}{\hat{s}}) + |C_9|^2[1 + 2\hat{s} + (1 - 4\hat{s})\frac{2t^2}{\hat{s}}] \\ &+ 12Re(C_8^{eff}C_7^*)(1 + \frac{2t^2}{\hat{s}}) + \frac{3}{2}|C_{Q_1}|^2(1 - \frac{4t^2}{\hat{s}})\hat{s} + \frac{3}{2}|C_{Q_2}|^2\hat{s} + 6Re(C_9C_{Q_2}^*)t \\ E(\hat{s}) &= Re(C_8^{eff}C_9^*\hat{s} + 2C_7C_9^* + C_8^{eff}C_{Q_1}^*t + 2C_7C_{Q_1}^*t) \end{aligned} \tag{6}$$

Another CP violating observable in $B \to X_s l^+ l^-$ is the normal polarization of the lepton in the decay, P_N, which is the T-violating projection of the lepton spin onto the normal of the decay plane, i.e $P_N \sim \vec{s}_l \cdot (\vec{p}_s \times \vec{p}_{l^-})$ [15]. A straightforward calculation leads to[8, 16]

$$P_N = \frac{3\pi}{4}\sqrt{1 - \frac{4t^2}{\hat{s}}}\hat{s}^{\frac{1}{2}} Im\Big[2C_8^{eff*}C_9t + 4C_9C_7^*\frac{t}{\hat{s}} + C_8^{eff*}C_{Q_1} + 2C_7^*C_{Q_1} + C_9^*C_{Q_2}\Big]\Big/D(\hat{s}) \tag{7}$$

The Wilson coefficients C_i and C_{Q_i} in eqs.(6) and (7) have been given in ref.[8, 17, 18]. Since only C_8^{eff} contains the non-trivial strong phase, A^1_{CP} is determined by ImC_7 and A^2_{CP} by ImC_{Q_1} and ImC_7. Although P_N depends on all the relevant Wilson coefficients a large P_N does require relatively large values of Im$C_{Q_i}(i = 1, 2)$ [8]. With the main contributions coming from exchanging chargino-stop loop with neutral Higgs coupled to external b quark [17], imaginary parts of C_{Q_i}s come mainly from terms proportional to (unitarity condition for stop mixing matrix has been used)

$$\frac{m_{\chi_i} m_t}{m_W^2 sin\beta\ cos\beta} U(i,2)V(i,2)D_{t21}D^*_{t11}, \qquad i=1,2 \tag{8}$$

i.e CP violating phases enter into the imaginary parts of C_{Q_i} through g-h mixings (U, V) and chiral mixing (D_t) of stops. From the chargino mass matrix

$$M_C = \begin{pmatrix} M_2 & \sqrt{2}m_W sin\beta \\ \sqrt{2}m_W cos\beta & \mu \end{pmatrix} \tag{9}$$

and stop mass matrix

$$M_{\tilde{t}}^2 = \begin{pmatrix} M_{\tilde{Q}}^2 + {m_t}^2 + M_z^2(\frac{1}{2} - Q_u \sin^2\theta_W)\cos 2\beta & m_t(A_t^* - \mu\cot\beta) \\ m_t(A_t - \mu^*\cot\beta) & M_{\tilde{U}}^2 + {m_t}^2 + M_z^2 Q_u \sin^2\theta_W \cos 2\beta \end{pmatrix}, \tag{10}$$

we know that $\sum_{i=1}^2 m_{\chi_i} U(i,2)V(i,2) = \mu$ and $D_{t21}D^*_{t11} = \frac{m_t}{m^2_{\tilde{t}_1} - m^2_{\tilde{t}_2}}(A_t - \mu^* cot\beta)$. Therefore, A_t itself is as important as μ for providing imaginary contributions to C_{Q_i}, in particular, for large tanβ. The similar conclusion holds also for C_7.

We have known well that A_t at the EW scale mainly depends on M_3 at the GUT scale through RGE effects[19, 20]. In fact there exists a quasi fixed point which shows the ratio of A_t at m_Z scale to M_3 at GUT scale to be about -1.6 provided the Yukawa couplings of the third generation are large enough [19]. Hence it is possible for A_t to achieve a very large imaginary parts in the non-universal gaugino mass models, in contrast to the case of mSUGRA. Especially in large $tan\beta$ case, $|\phi_\mu|$ is limited by EDM data to be less than $\pi/6$, so A_t plays a more important role in CP Violation than μ.

To study the effect of large ϕ_3 (hence large ϕ_{A_t}) on C_7, we notice that essentially A_t is multiplied by μ. Changing the sign of A_t has the same effects of switching the sign of μ and switching the sign of μ results in a sign change in C_7 (because in most of the parameter space μ are much larger than the non-diagonal terms in eq.(9)), so if ϕ_3 is in $[\pi/2, 3\pi/2]$ (hence ϕ_{A_0} in $[-\pi/2, \pi/2]$) and ϕ_μ in $[-\pi/2, \pi/2]$), supersymmetry contributions give enhancement to ReC_7 so that it is hard to satisfy the $B \to X_s\gamma$ constraints: $2.\times 10^{-4} < Br(B \to X_s\gamma) < 4.5 \times 10^{-4}$ [21]. Similar situation occurs for ϕ_3 being in $[-\pi/2, \pi/2]$) and ϕ_μ in $[\pi/2, 3\pi/2]$. Since supersymmetry gives large contributions to C_7 only when $tan\beta$ is large we shall focus on large $tan\beta$ case in the following. From the above analysis of EDMN we know that for somewhat large ϕ_μ(say around ∓ 0.4) cancellations happen for ϕ_3 near $\pm\pi/2$. With this kind of phases of M_3 and μ SUSY contributions to C_7 are almost totally imaginary. So the value of ImC_7 is constrained to be very small by the branching ratio of $B \to X_s\gamma$. One way to avoid the $B \to X_s\gamma$ constraint is to make use of the cancellation happened at ϕ_3 around $\pm\pi/4$ and ϕ_μ about ∓ 0.2, With such kind of phases and a low mass spectrum (M_2 and M_3 around 150 Gev) the real part of SUSY contributions to C_7 cancels those from W-top and charged Higgs-top loops. The real part can even be cancelled to be near zero and only a large imaginary part of C_7 remains. Another way is to just suppress the total SUSY contributions to C_7, i.e., to make the mass spectrum heavier but still less than 1 Tev (M_2 and M_3 larger than about 300 Gev and less than about 500 Gev). In the former case A^1_{CP} can reach order 1% for $B \to X_s e^+e^-$. For $B \to X_s\mu^+\mu^-$ and $B \to X_s\tau^+\tau^-$, because of their larger Yukawa coupling there are great enhancement of branching ratio[17] so that A^1_{CP} are smaller than that for $B \to X_s e^+ e^-$. In the later case(we shall call it as region A hereafter) it can only be a few thousandth at most, i.e., the same order as that in SM.

As pointed out above, the effects of large ϕ_3 on C_{Q_i} are similar to those on C_7. In large tanβ case, for $\phi_\mu \sim \pm 0.4$ and $\phi_3 \sim \mp\pi/2$ or $\phi_\mu \sim \pi \pm 0.4$ and $\phi_3 \sim \pm\pi/2$ (which we shall call as region B for simplisity), which are allowed by the EDME and EDMN limits, ImC_{Q_i} reachs maxima. In small tanβ case, although the constraints of EDMs on ϕ_μ and ϕ_3 are relaxed the magnitude of ImC_{Q_i} is very small since C_{Q_i} is proportional to $m_l tan^2\beta$(even $m_l tan^3\beta$ in some regions of the parameter space). Therefore, we expect the significant CP violation in large tanβ case.

From eqs.(5) and (6), A^2_{CP} can be rewritten as

$$A^2_{CP} = \frac{E(\hat{s})\overline{D}(\hat{s}) - \overline{E}(\hat{s})D(\hat{s})}{E(\hat{s})\overline{D}(\hat{s}) + \overline{E}(\hat{s})D(\hat{s})}$$

For l=e, the difference between $E(\hat{s})$ and $\overline{E}(\hat{s})$ can be neglected(as it is proportional to lepton mass square, see eq.(6)). So A^2_{CP} for l=e can be reduced to

$$A^2_{CP} \doteq \frac{\overline{D}(\hat{s}) - D(\hat{s})}{\overline{D}(\hat{s}) + D(\hat{s})}$$

that is exactly the opposite of A^1_{CP}. The same conclusion can be drawn for l=μ, τ in small tanβ case due to smallness of C_{Q_1}. On the other hand, for l=τ in large tanβ case, $|E(\hat{s}) - \overline{E}(\hat{s})|$ can be more important than $|D - \bar{D}|$ and consequently one has approximately

$$A^2_{CP} \doteq \frac{E(\hat{s}) - \overline{E}(\hat{s})}{E(\hat{s}) + \overline{E}(\hat{s})}$$

Thus, it is propotional to t ImC_{Q1}. Therefore, in region B where C_{Q_i}s reach the maxima, A^2_{CP} can be over 50%. In region A, C_{Q_i}s are less important and A^2_{CP} can reach about 5% at most. The correlation between A^2_{CP} and EDME (or EDMN) in region A is plotted in fig.3 (note that we choose $\hat{s} = 0.76$ as representative in the figure). For l=μ, the magnitude of A^2_{CP} can be estimated to be of order 1% at most in the large tanβ case. Numerical calculations prove this estimate.

Fig.4 shows the correlation of EDM constraints and P_N of $B \to X_s \tau^+ \tau^-$in region A. We can see in this figure that P_N can reach more than 15 percent. In region B, as C_{Q_i}s are much larger the numerator in eq.(7) is increased a lot. But the denominator in eq.(7) (and consenquently the branching ratio of $B \to X_s \tau^+ \tau^-$) is also greatly enhanced in this region, so P_N is just about 15%, i.e., not larger than the magnitude that can be achieved in region A. Situations for muon are similar and because of its much smaller Yukawa coupling the magnitude of P_N can only reach about 6%. For electron P_N is negligibly small, due to its neglegibly small mass. An important feature that can be seen from fig.3 and fig.4 is that the magnitudes of A^2_{CP} and P_N will not be reduced if EDM constraints improved. That is because the regions of parameter space where EDM constraints are satisfied are of width about $\pi/20$ for ϕ_3 and about $\pi/4$ for ϕ_{A_0}(adjustment needed), while C_{Q_i}s do not change sharply within these regions.

In summary, we have analyzed the constraints of electric dipole moments of electron and neutron on the parameter space in supergravity models with nonuniversal gaugino masses. It is shown that with a light sparticle spectrum, the sufficient cancellations in the calculation of EDMs can happen due to the presence of the two new phases arising from complex gaugino masses, in addition to the phases ϕ_μ and ϕ_{A_0}. With appropriate correlation between ϕ_μ and ϕ_1 (for EDME) or ϕ_3 (for EDMN) as well as an appropriate choice of ϕ_{A_0}, cancellations can occur and all phases can be order of one in the small tanβ case and all phases but ϕ_μ ($|\phi_\mu| \lesssim \pi/6$) order of one in the large tanβ case. This is in contrast to the case of mSUGRA where in the parameter space where cancellations among various SUSY contributions to EDMs happen ϕ_μ must be less than $\pi/10$ for small $tan\beta$ and $\mathcal{O}(10^{-2})$ for large $tan\beta$. And our analysis show that the branching ratio of $B \to X_s\gamma$ gives an extra constraint on the phases for large tanβ case with light mass spectrum. We have calculated direct CP asymmetries and the CP violating normal polarization of lepton in $B \to X_s l^+ l^-$ in the regions of the parameter space in the models where the constraints from EDMs as well as $B \to X_s\gamma$ are satisfied. It is shown that the results for A^1_{CP} are similar to those in mSUGRA if the mass sprectrum is relatively heavier($M_2 \& M_3 \gtrsim 300 Gev$) and it is also true for A^2_{CP} in the small tanβ case. The former is due to the constraint from $B \to X_s\gamma$ and the latter is due to smallness of the contributions from exchanging neutral Higgs bosons in the small tanβ case. However, in the large tanβ case, A^2_{CP} can reach 1% for $l = \mu$ and is a few percent in most of allowed regions and can reach 50% in some allowed regions for l=τ. A^2_{CP} for $l = e$ is approximately equal to A^1_{CP} even in the large tanβ case. P_N can reach 6% for $l = \mu$ and is in the range from 1% to 15% in most of the allowed regions for l=τ in the large tanβ case. That is, there is a significant enhancement compared to the mSUGRA in which P_N only can reach about 0.5% for l=μ and about 5% for l=τ. In the small tanβ case the results are similar to those in mSUGRA. For l=e, P_N is negligibly small, as it should be.

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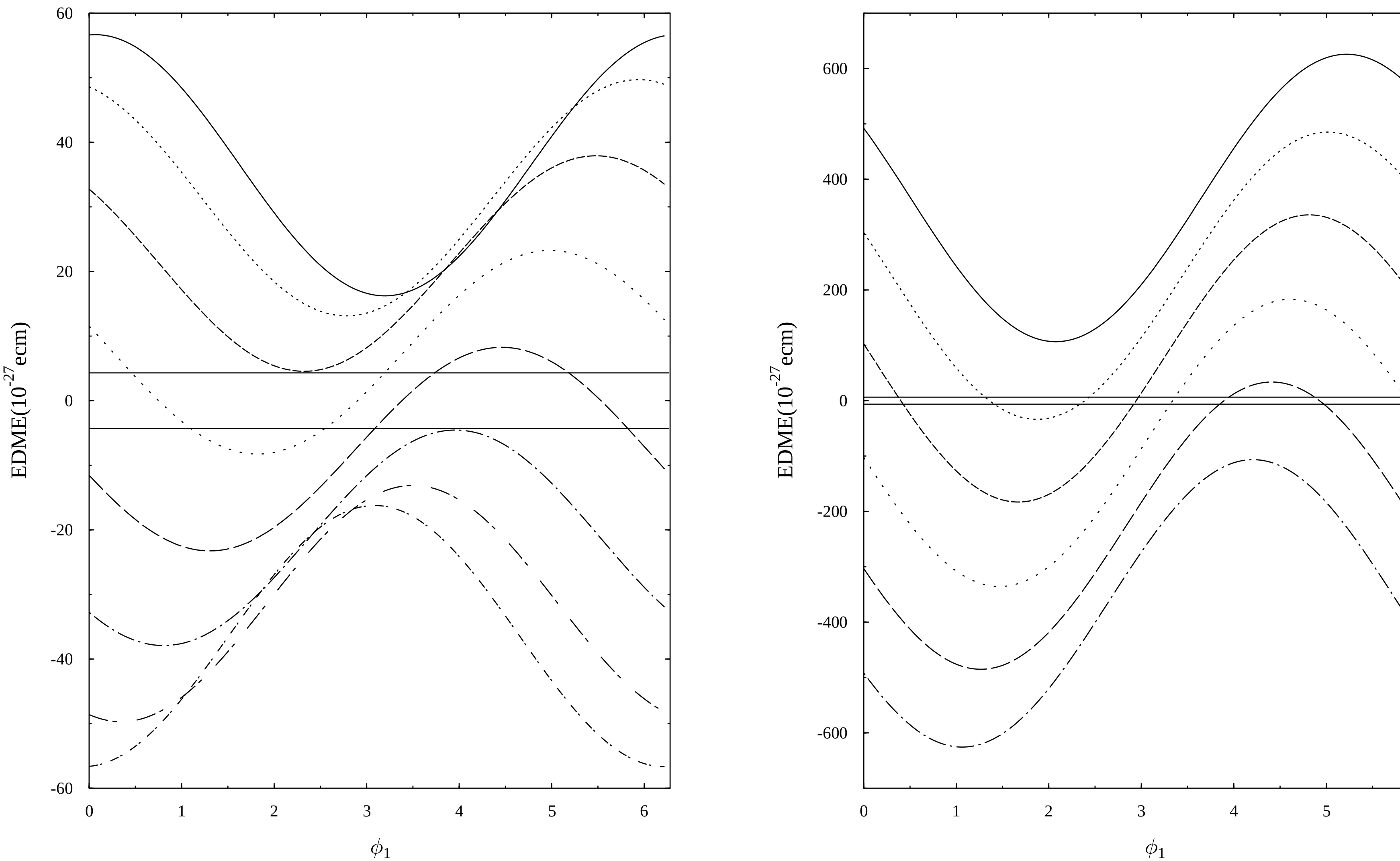

Fig1. EDME(10^{-27}) as functions of ϕ_1 from 0 to 2π. (a) $\tan\beta$=2 and lines from below refer to ϕ_μ=1.4, 1.0, 0.6, 0.2, -0.2, 0.6 ,1.4.(b) $\tan\beta$=30 and six lines from below refer to ϕ_μ=0.5, 0.3 0.1, -0.1, 0.3, 0.5. Phases of M_2, M_3 and A_0 are chosen to be Other parameters are chosen such that M_0=$|M_2|$=400Gev, $|M_1|$=$|M_3|$=500Gev, $|A_0|$=800Gev.

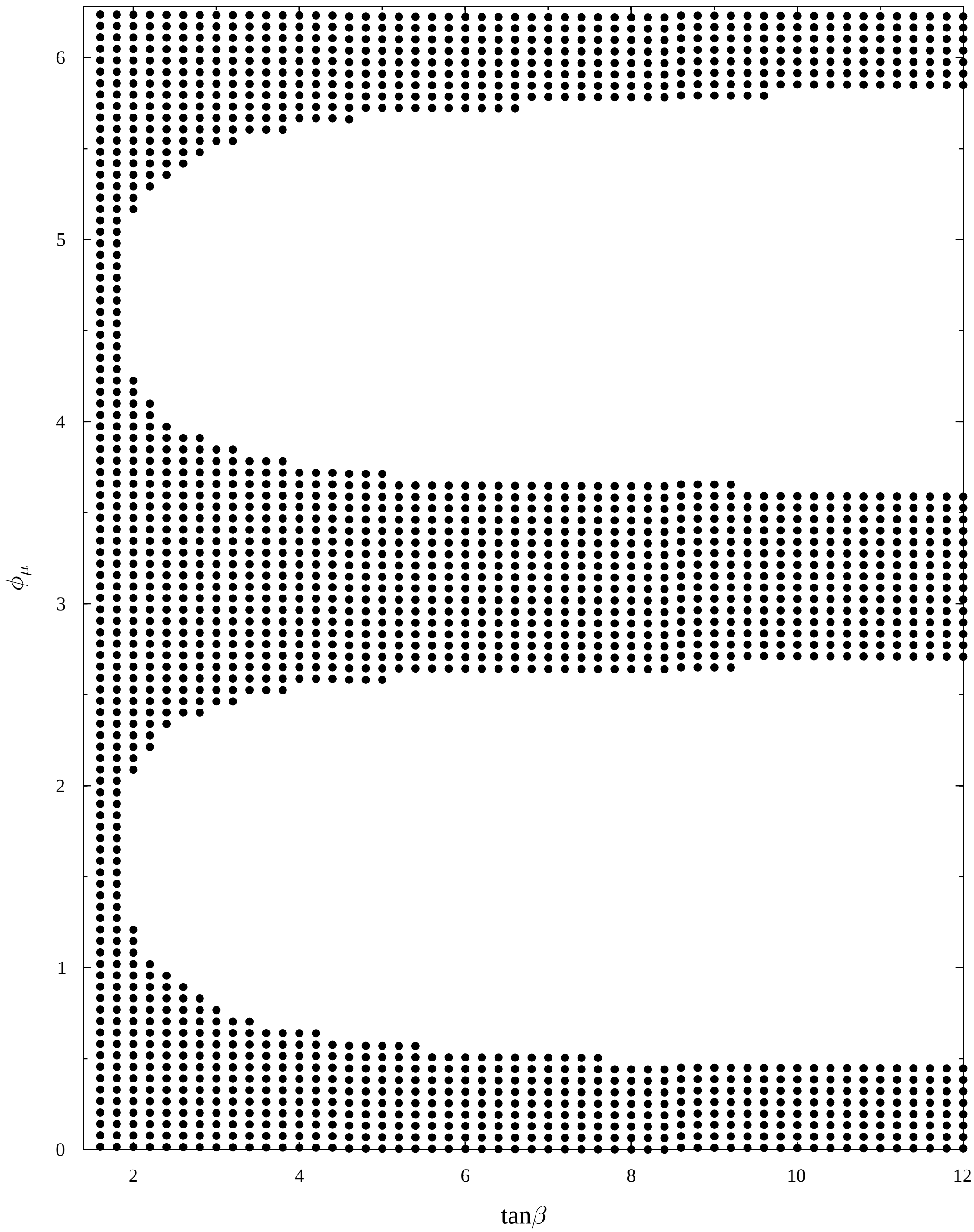

Fig.2 Correlation between the allowed regions of ϕ_μ and $\tan\beta$ with $M_0=|M_2|=400$Gev, $|M_1|=|M_3|=500$Gev, $|A_0|=800$Gev.

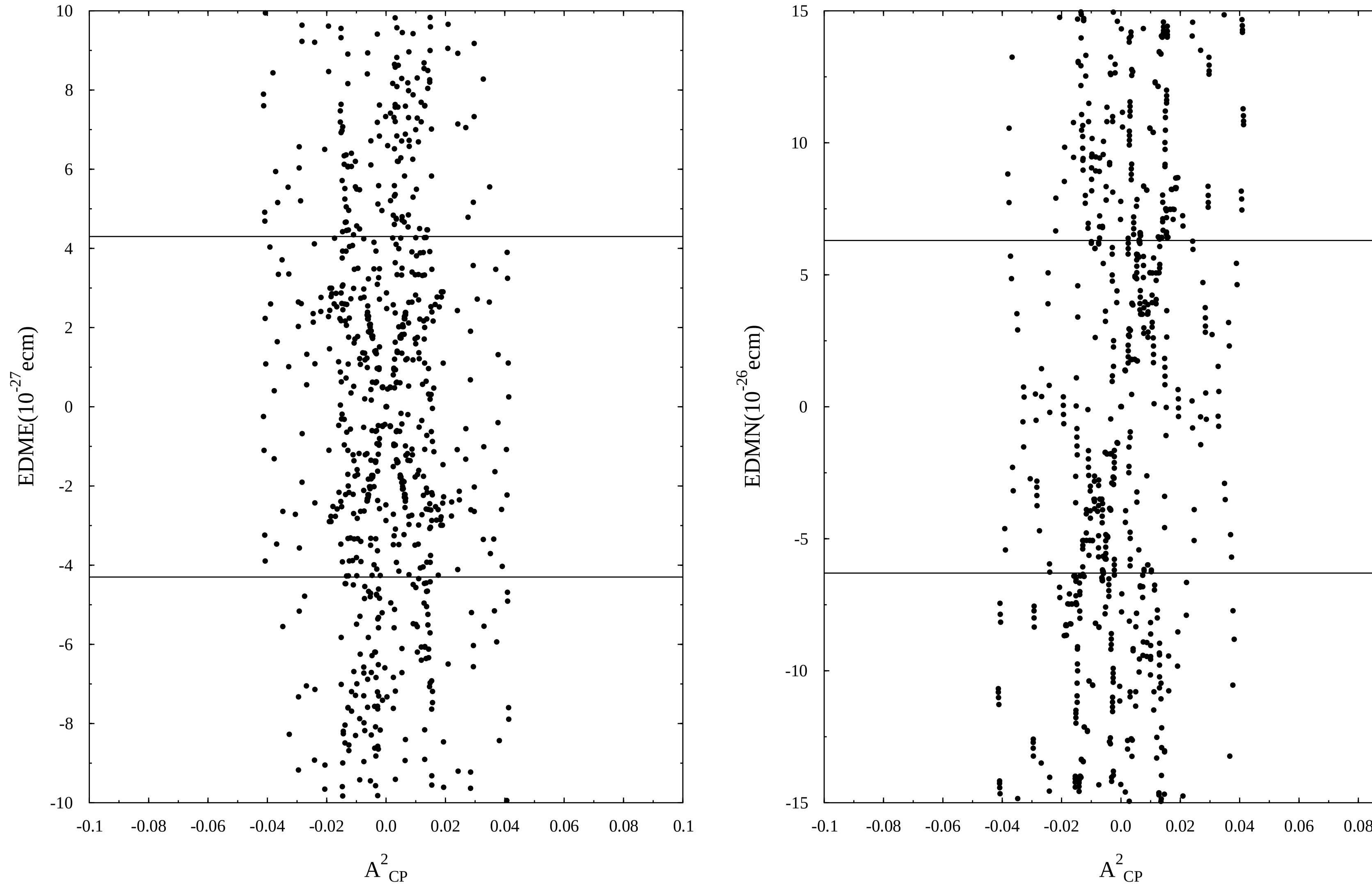

Fig.3 Correlation of EDMs and A^2_{CP} in region A.

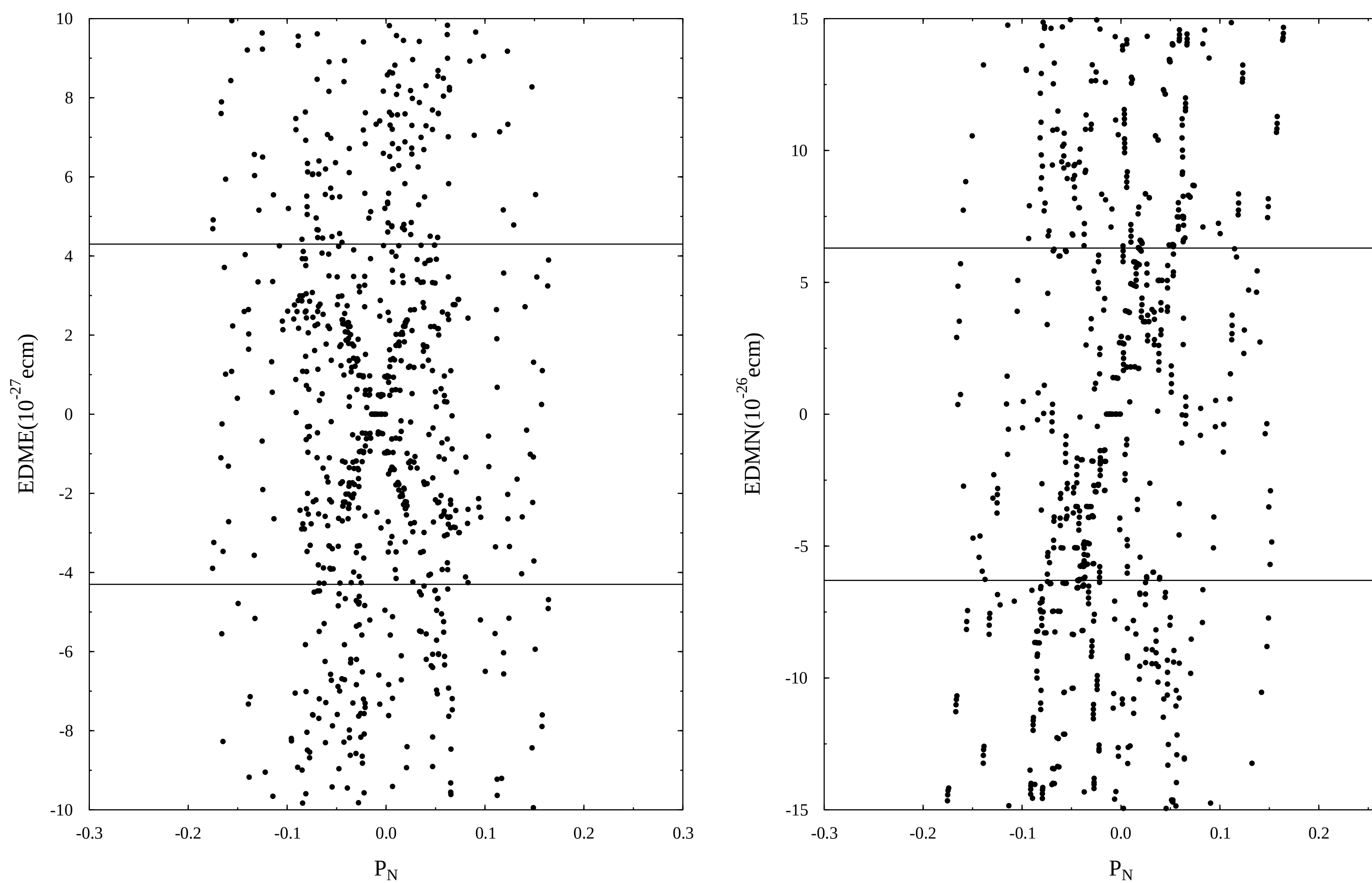

Fig.4 Correlation of EDMs and P_N in region A.